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Authors
Cho, Peter
Simmons, Elizabeth H

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Looking for Gluon Substructure at the Tevatron

Peter Cho†
Lauritsen Laboratory
California Institute of Technology
Pasadena, CA 91125

and

Elizabeth Simmons‡
Lyman Laboratory of Physics
Harvard University
Cambridge, MA 02138

Abstract

The impact of nonrenormalizable gluon operators upon inclusive jet cross sections is studied. Such operators could arise in an effective strong interaction Lagrangian from gluon substructure and would induce observable cross section deviations from pure QCD at high transverse jet energies. Comparison of the theoretical predictions with recent CDF data yields a lower limit on the gluon compositeness scale Λ. We find $\Lambda > 2.03$ TeV at 95% CL.

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The inclusive jet cross section data from the 1988-89 Fermilab Tevatron run span seven orders of magnitude and include the highest transverse jet energy measurements reported to date [1]. These data provide a stringent test of quantum chromodynamics and constrain possible new physics beyond the Standard Model. In particular, they set improved limits on hypothetical quark substructure. At energies small compared to the compositeness scale Λ, the dominant effects from quark substructure can be reproduced by four-quark operators in a low energy effective Lagrangian [2]. Although their coefficients are unknown in the absence of a detailed theory of preon dynamics, the general impact of these nonrenormalizable operators upon parton scattering may be estimated. The remarkable agreement between the experimental measurements and the predictions of QCD then places a bound on the quark compositeness scale. CDF finds a lower limit on Λ of 1.4 TeV at the 95% confidence level [1].

In this letter, we reinterpret the CDF data to probe for signals of new physics that could arise in the gluon sector. Specifically, we consider the impact upon the inclusive cross section measurements of nonrenormalizable gluonic operators which may appear in the effective strong interaction Lagrangian. Such operators could originate from a number of different sources. For example, suppose there exist new heavy colored bosons or fermions beyond those in the Standard Model. Such particles would induce nonlocal interactions among gluons through loop diagrams. The leading behavior of these graphs can readily be extracted and reexpressed via an operator product expansion in terms of local but nonrenormalizable gluon operators. Alternatively, we might speculate that gluons are bound states of some more fundamental preon constituents. Then as in the case of composite quarks, preon exchange could generate nonrenormalizable gluon interactions. In the following, we will adopt a model independent approach and not specify the underlying physics whose low energy effects are encoded in the effective Lagrangian. Instead, we simply seek to place a limit on its characteristic scale Λ.

We first enumerate the lowest dimension gluon operators whose scattering effects would be easiest to observe. There exist only two independent operators of mass dimension $d+2$ in $d=4-\epsilon$ spacetime dimensions which preserve gauge invariance along with C, P and T [3]:

$$\begin{aligned} O_1 &= \frac{\mu^{\epsilon/2} g}{\Lambda^2} f_{abc} G^{\rho}_{a\nu} G^{\nu}_{b\lambda} G^{\lambda}_{c\rho} \\ O_2 &= \frac{1}{2!\Lambda^2} D^{\rho} G^{a}_{\rho\nu} D_{\lambda} G^{\lambda\nu}_{a} \end{aligned} \tag{1}$$

where

$$
\begin{aligned}
D^\rho &= \partial^\rho - i\mu^{\epsilon/2} g G_a^\rho T_a \\
G_a^{\rho\nu} &= \partial^\rho G_a^\nu - \partial^\nu G_a^\rho + \mu^{\epsilon/2} g f_{abc} G_b^\rho G_c^\nu
\end{aligned}
\tag{2}
$$

and μ denotes the renormalization scale. All other dimension-$(d+2)$ gluon operators either vanish or reduce to combinations of the two in eqn. (1). [1]

The contributions of O_1 and O_2 to parton scattering cross sections have been studied in refs. [3,4]. Operator O_1 mediates gluon-quark as well as gluon-gluon scattering and interferes with the pure QCD amplitudes for these same processes. One might expect its impact to become pronounced at higher energies due to the sizable gluon content of colliding hadrons at small parton momentum fractions. However, the helicity structure of the $O(1/\Lambda^2)$ amplitude from O_1 for $gg \to gg$ scattering is orthogonal to that of pure QCD. So the two amplitudes do not interfere. Similarly, the $O(1/\Lambda^2)$ interference terms in the differential cross sections for $gg \to q\overline{q}$ and the other partonic processes related by crossing vanish in the limit of zero quark mass. Thus the presence of O_1 in the low energy Lagrangian affects inclusive jet cross sections starting only at $O(1/\Lambda^4)$. This surprising and rather disappointing result motivates us to search for signals of gluon compositeness in non-gluonic channels.

Fortunately, operator O_2 mediates quark-quark scattering. The classical equation of motion

$$
D_\rho G_a^{\rho\nu} = -\mu^{\epsilon/2} g \sum_{\text{flavors}} \overline{q}\gamma^\nu T_a q \tag{3}
$$

relates the S-matrix elements of O_2 to those of a color octet four-quark operator:

$$
O_2 \xrightarrow{EOM} \frac{\mu^\epsilon g^2}{2!\Lambda^2} \sum_{\text{flavors}} \big(\overline{q}\gamma_\nu T_a q\big)\big(\overline{q}\gamma^\nu T_a q\big). \tag{4}
$$

This identification should be understood as a relation among S-matrix elements and not as a true operator identity [5]. Notice that the equation of motion (3) produces a factor of g^2 multiplying the operator in (4). No such extra powers of the strong interaction coupling accompany the analogous color singlet operators that enter into quark compositeness analyses. At typical Tevatron energies, the numerical value for g^2 is somewhat larger than

[1] The operator obtained from O_1 by replacing the antisymmetric structure constant f_{abc} with the completely symmetric symbol d_{abc} violates charge conjugation and is identically zero.

unity. So we anticipate that the effect upon quark scattering from the four-quark operator induced by gluon substructure will be slightly larger than that from the four-quark operators generated by quark substructure.

To complete our operator basis so that it closes under renormalization, we include three more four-quark operators along with O_1 and O_2 in the effective Lagrangian:

$$\mathcal{L}_{\rm eff} = \mathcal{L}_{QCD} + \sum_{i=1}^{5} C_i(\mu) O_i(\mu) \tag{5}$$

where

$$\begin{aligned}
O_1 &= \frac{\mu^{\epsilon/2} g}{\Lambda^2} f_{abc} G^{\rho}_{a\nu} G^{\nu}_{b\lambda} G^{\lambda}_{c\rho} \\
O_2 &= \frac{\mu^{\epsilon} g^2}{2!\Lambda^2} \sum_{\rm flavors} \big(\overline{q}\gamma_\nu T_a q\big)\big(\overline{q}\gamma^\nu T_a q\big) \\
O_3 &= \frac{\mu^{\epsilon} g^2}{2!\Lambda^2} \sum_{\rm flavors} \big(\overline{q}\gamma_\nu \gamma^5 T_a q\big)\big(\overline{q}\gamma^\nu \gamma^5 T_a q\big) \\
O_4 &= \frac{\mu^{\epsilon} g^2}{2!\Lambda^2} \sum_{\rm flavors} \big(\overline{q}\gamma_\nu q\big)\big(\overline{q}\gamma^\nu q\big) \\
O_5 &= \frac{\mu^{\epsilon} g^2}{2!\Lambda^2} \sum_{\rm flavors} \big(\overline{q}\gamma_\nu \gamma^5 q\big)\big(\overline{q}\gamma^\nu \gamma^5 q\big).
\end{aligned} \tag{6}$$

Since the underlying short-distance physics responsible for generating these nonrenormalizable operators is not known, there is a fair amount of arbitrariness in the values one chooses for their dimensionless coefficients and the compositeness scale Λ. We will follow the convention adopted in previous quark substructure studies and define Λ to be the scale where the magnitude of the gluon operator coefficients C_1 or C_2 equals 4π [2,6]. We set the coefficients of the remaining operators at the Λ scale to zero.

To determine the operators' coefficients at energies probed by the Tevatron, we evolve their values down from the compositeness scale via the renormalization group equation

$$\mu \frac{d}{d\mu} C_i(\mu) = \sum_j (\gamma^T)_{ij} C_j(\mu). \tag{7}$$

The anomalous dimension matrix

$$\gamma = \begin{array}{c} \\ O_1 \\ O_2 \\ O_3 \\ O_4 \\ O_5 \end{array} \!\!\begin{array}{c} \begin{array}{ccccc} O_1 & O_2 & O_3 & O_4 & O_5 \end{array} \\ \begin{pmatrix} 7+2n_f/3 & 0 & 0 & 0 & 0 \\ 0 & 311/36-2n_f/3 & 5/4 & 0 & 2/3 \\ 0 & 41/36 & 35/4-2n_f/3 & 2/3 & 0 \\ 0 & 4/3 & 6 & 11-2n_f/3 & 0 \\ 0 & 22/3 & 0 & 0 & 11-2n_f/3 \end{pmatrix} \end{array} \frac{g^2}{8\pi^2}$$

describes the mixing for n_f active quark flavors among all the dimension-$(d+2)$ operators in the effective Lagrangian. The last four rows of γ quantify the running of the four-quark operators in our basis and were determined from a straightforward operator mixing computation. The entries in the first row on the other hand were extracted from the highly nontrivial anomalous dimension calculations reported in ref. [7] by Narison and Tarrach and in ref. [8] by Morozov. Notice that the triple gluon field strength operator runs only into itself and does not mix with any of the other four-quark operators. [2] Since O_1 has virtually no impact upon parton scattering and does not mix into operators whose effects can be detected, it is essentially invisible at $O(1/\Lambda^2)$. We will consequently be unable to place any limit upon its associated scale.

It is convenient to decompose the anomalous dimension matrix as

$$\begin{aligned}\gamma &= \frac{g^2}{8\pi^2}\hat{\gamma} + O(g^4) \\ &= \frac{g^2}{8\pi^2} SDS^{-1} + O(g^4)\end{aligned} \tag{8}$$

where the eigenvalues λ_i of $\hat{\gamma}$ are contained in the diagonal matrix D while the corresponding eigenvectors are arranged into the columns of matrix S. The general solution to the coefficients' renormalization group equation

$$C_i(\mu) = \sum_j \Big[\exp \int_{g(\Lambda)}^{g(\mu)} dg \frac{\gamma^T(g)}{\beta(g)}\Big]_{ij} C_j(\Lambda) \tag{9}$$

can then be rewritten as

$$C_i(\mu) = \sum_{j,k} (S^{-1})^T_{ij} \mathrm{diag}\Big(\Big[\frac{\alpha_s(\mu)}{\alpha_s(\Lambda)}\Big]^{\frac{\lambda_j}{2b}}\Big) (S^T)_{jk}\, C_k(\Lambda) \tag{10}$$

where $b = -11/2 + n_f/3$ is the coefficient in the one-loop QCD beta function $\beta(g) = bg^3/8\pi^2$ [10]. Integrating β, we obtain the strong interaction fine structure constant

$$\alpha_s(\mu) = \frac{\alpha_s(M_z)}{1 - b\dfrac{\alpha_s(M_z)}{\pi}\log\dfrac{\mu}{M_z}}. \tag{11}$$

[2] As shown in Morozov's paper, the anomalous dimension of the CP-even G^3 operator is the same as that of the CP-odd $\tilde{G}^3$ operator which gained much notoriety a few years ago [9].

We choose the constant of integration in the integrated beta function to be $\alpha_s(M_z) = 0.118 \pm 0.007$ [11] rather than the QCD scale since parton energies at the Tevatron are typically large compared to the Z scale.

We now turn to computing the inclusive jet cross section in $p\overline{p}$ collisions. Neglecting higher order multi-jet events, we start with the two-jet differential cross section

$$\frac{d^3\sigma}{d\eta_1 d\eta_2 dp_T}(AB \to 2\text{jets}) = \sum_{abcd} 2x_a x_b p_T \Big[f_{a/A}(x_a) f_{b/B}(x_b) + (A \leftrightarrow B \text{ if } a \neq b)\Big] \frac{d\sigma}{d\hat{t}}(ab \to cd) \tag{12}$$

expressed in terms of the jets' pseudorapidities (η_1, η_2), their common transverse momentum (p_T), and the momentum fractions (x_a, x_b) of partons a and b inside hadrons A and B. The partons' distribution functions $f_{a/A}(x_a)$ and $f_{b/B}(x_b)$ are folded together with the differential cross section $d\sigma/d\hat{t}$ for the elementary scattering process $ab \to cd$. [3] The product is then summed over all possible initial and final parton configurations. To convert the two-jet expression (12) into an inclusive single-jet cross section, we integrate over the pseudorapidity range of one jet, average the other over the pseudorapidity interval $0.1 \leq |\eta| \leq 0.7$ visible to the CDF detector, and multiply by two to count the contributions of both jets to the inclusive cross section:

$$\frac{1}{\Delta\eta} \int d\eta \frac{d^2\sigma}{d\eta dp_T} = \frac{2}{\Delta\eta} \int d\eta_1 \int_{-\infty}^{\infty} d\eta_2 \frac{d^3\sigma}{d\eta_1 d\eta_2 dp_T}. \tag{13}$$

The result may then be compared with the measurements reported by CDF [1].

The lowest order QCD predictions for the parton cross sections

$$\frac{d\sigma}{d\hat{t}}(ab \to cd) = \frac{\pi\alpha_s^2}{\hat{s}^2} \Sigma(ab \to cd) \tag{14}$$

have been frequently documented in the literature [12,13]. They conventionally include initial state color averaging factors and are written in terms of the partonic invariants $\hat{s}$, $\hat{t}$ and $\hat{u}$. The QCD formulae for $\Sigma(ab \to cd)$ are modified by the nonrenormalizable operators in our effective Lagrangian which induce the $O(1/\Lambda^2)$ interference terms tabulated below:

[3] Note that the Mandelstam invariant $\hat{t}$ refers to the colliding partons rather than the incident hadrons.

$$\Sigma(qq'\to qq') = \frac{4}{9}\frac{\hat{s}^2+\hat{u}^2}{\hat{t}^2} + \frac{8}{9}\frac{(C_2+C_3)\hat{s}^2+(C_2-C_3)\hat{u}^2}{\hat{t}\Lambda^2} + O\Big(\frac{1}{\Lambda^4}\Big) \tag{15a}$$

$$\Sigma(q\overline{q}\to q'\overline{q}') = \frac{4}{9}\frac{\hat{t}^2+\hat{u}^2}{\hat{s}^2} + \frac{8}{9}\frac{(C_2+C_3)\hat{u}^2+(C_2-C_3)\hat{t}^2}{\hat{s}\Lambda^2} + O\Big(\frac{1}{\Lambda^4}\Big) \tag{15b}$$

$$\begin{aligned}\Sigma(qq\to qq) = {} & \frac{4}{9}\Big(\frac{\hat{s}^2+\hat{u}^2}{\hat{t}^2} + \frac{\hat{s}^2+\hat{t}^2}{\hat{u}^2}\Big) - \frac{8}{27}\frac{\hat{s}^2}{\hat{t}\hat{u}} \\ & + \frac{8C_2}{9\Lambda^2}\Big(\frac{\hat{s}^2+\hat{u}^2}{\hat{t}} + \frac{\hat{s}^2+\hat{t}^2}{\hat{u}}\Big) + \frac{8C_3}{9\Lambda^2}\Big(\frac{\hat{s}^2-\hat{u}^2}{\hat{t}} + \frac{\hat{s}^2-\hat{t}^2}{\hat{u}}\Big) \\ & + \Big(\frac{8(C_2+C_3)}{27\Lambda^2} - \frac{16(C_4+C_5)}{9\Lambda^2}\Big)\frac{\hat{s}^3}{\hat{t}\hat{u}} + O\Big(\frac{1}{\Lambda^4}\Big)\end{aligned} \tag{15c}$$

$$\begin{aligned}\Sigma(q\overline{q}\to q\overline{q}) = {} & \frac{4}{9}\Big(\frac{\hat{s}^2+\hat{u}^2}{\hat{t}^2} + \frac{\hat{t}^2+\hat{u}^2}{\hat{s}^2}\Big) - \frac{8}{27}\frac{\hat{u}^2}{\hat{s}\hat{t}} \\ & + \frac{8C_2}{9\Lambda^2}\Big(\frac{\hat{s}^2+\hat{u}^2}{\hat{t}} + \frac{\hat{t}^2+\hat{u}^2}{\hat{s}}\Big) - \frac{8C_3}{9\Lambda^2}\Big(\frac{\hat{s}^2-\hat{u}^2}{\hat{t}} + \frac{\hat{t}^2-\hat{u}^2}{\hat{s}}\Big) \\ & + \Big(\frac{8(C_2+C_3)}{27\Lambda^2} - \frac{16(C_4+C_5)}{9\Lambda^2}\Big)\frac{\hat{u}^3}{\hat{s}\hat{t}} + O\Big(\frac{1}{\Lambda^4}\Big)\end{aligned} \tag{15d}$$

$$\Sigma(gg\to q\overline{q}) = \frac{1}{6}\Big(\frac{\hat{t}^2+\hat{u}^2}{\hat{t}\hat{u}}\Big) - \frac{3}{8}\frac{\hat{t}^2+\hat{u}^2}{\hat{s}^2} + O\Big(\frac{1}{\Lambda^4}\Big) \tag{15e}$$

$$\Sigma(q\overline{q}\to gg) = \frac{32}{27}\Big(\frac{\hat{t}^2+\hat{u}^2}{\hat{t}\hat{u}}\Big) - \frac{8}{3}\frac{\hat{t}^2+\hat{u}^2}{\hat{s}^2} + O\Big(\frac{1}{\Lambda^4}\Big) \tag{15f}$$

$$\Sigma(gq\to gq) = -\frac{4}{9}\Big(\frac{\hat{s}^2+\hat{u}^2}{\hat{s}\hat{u}}\Big) + \frac{\hat{s}^2+\hat{u}^2}{\hat{t}^2} + O\Big(\frac{1}{\Lambda^4}\Big) \tag{15g}$$

$$\Sigma(gg\to gg) = \frac{9}{2}\Big(3 - \frac{\hat{t}\hat{u}}{\hat{s}^2} - \frac{\hat{s}\hat{u}}{\hat{t}^2} - \frac{\hat{s}\hat{t}}{\hat{u}^2}\Big) + O\Big(\frac{1}{\Lambda^4}\Big). \tag{15h}$$

Here q' denotes a quark not identical in flavor to quark q. We have dropped $O(1/\Lambda^4)$ terms in these formulae since we are not keeping track of any dimension-$(d+4)$ gluon operators whose contributions to $\Sigma(ab\to cd)$ are of the same order. We have also neglected all parton masses. At transverse jet energies of a few hundred GeV, this should be a good approximation except for the top quark whose mass we assume is $m_t = 140$ GeV. Initial state top mass effects may be safely ignored as the top content of colliding protons and antiprotons is negligible. However for processes involving $t\overline{t}$ production, we replace eqns. (15b) and (15e) with the heavy flavor QCD cross sections given in ref. [14] and incorporate $O(1/\Lambda^2)$ interference corrections:

$$\Sigma(q\bar{q}\to t\bar{t}) = \frac{4}{9}\frac{(\hat{t}-m_t^2)^2+(\hat{u}-m_t^2)^2+2m_t^2\hat{s}}{\hat{s}^2} + \frac{8}{9}\frac{(C_2+C_3)\hat{u}^2+(C_2-C_3)\hat{t}^2+2C_2m_t^2\hat{s}}{\hat{s}\Lambda^2} + O\Big(\frac{1}{\Lambda^4}\Big) \tag{16a}$$

$$\begin{aligned}\Sigma(gg\to t\bar{t}) = {} & \frac{3}{4}\frac{(m_t^2-\hat{t})(m_t^2-\hat{u})}{\hat{s}^2} - \frac{1}{24}\frac{m_t^2(\hat{s}-4m_t^2)}{(m_t^2-\hat{t})(m_t^2-\hat{u})} \\ & + \frac{1}{6}\frac{(m_t^2-\hat{t})(m_t^2-\hat{u})-2m_t^2(m_t^2+\hat{t})}{(m_t^2-\hat{t})^2} + \frac{1}{6}\frac{(m_t^2-\hat{t})(m_t^2-\hat{u})-2m_t^2(m_t^2+\hat{u})}{(m_t^2-\hat{u})^2} \\ & - \frac{3}{8}\frac{(m_t^2-\hat{t})(m_t^2-\hat{u})+m_t^2(\hat{u}-\hat{t})}{\hat{s}(m_t^2-\hat{t})} - \frac{3}{8}\frac{(m_t^2-\hat{t})(m_t^2-\hat{u})+m_t^2(\hat{t}-\hat{u})}{\hat{s}(m_t^2-\hat{u})} \\ & - \frac{9}{8}\frac{C_1m_t^2}{\Lambda^2}\Big(3+2m_t^2\frac{\hat{t}^2+\hat{t}\hat{u}+\hat{u}^2}{\hat{s}\hat{t}\hat{u}} - 3m_t^4\frac{\hat{t}+\hat{u}}{\hat{s}\hat{t}\hat{u}}\Big) + O\Big(\frac{1}{\Lambda^4}\Big).\end{aligned} \tag{16b}$$

Note that coefficient C_1 of the triple gluon field strength operator enters into eqn. (16b). But since it does not appear in any other cross section formula at $O(1/\Lambda^2)$, operator O_1 has almost no perceptible effect upon the inclusive cross section. We consequently ignore it from here on.

Combining eqns. (12) – (16), we calculate the single-jet inclusive cross section as a function of transverse jet energy E_T. We perform the computation using the leading order parton distribution functions of Morfin and Tung (MT set SL) [15] and the CTEQ collaboration (CTEQ set L) [16] as well as the next-to-leading order functions of Morfin and Tung (MT sets B1 and S), Harriman, Martin, Roberts and Stirling (HMRS set B) [17] and the CTEQ collaboration (CTEQ set MS). All of these parton distribution functions along with several others are conveniently contained within the PAKPDF package [18].

Representative results obtained from the MT set SL structure function evaluated at the renormalization scale $Q^2 = E_T^2/2$ are compared with the experimental data in fig. 1. The solid curve in the figure illustrates the predictions of pure QCD with no nonrenormalizable operator interactions. Following the example of the CDF analysis [1], we have multiplied the theoretical predictions by a normalization factor n to align them with the data. A fit for this constant performed over the region 80 GeV $\le E_T \le$ 160 GeV where effects from any compositeness operator terms are negligible yields $n = 1.35 \pm 0.01$. The resulting agreement between the shapes of the QCD and experimental cross section values is striking. There is however a slight suggestion of discrepancy at the highest measured transverse energies where compositeness operator effects would be expected to first show up. We therefore plot in the same figure the differential cross sections obtained after setting $\Lambda = 2.0$ TeV and $C_2(\Lambda) = -4\pi$ in our effective Lagrangian. The resulting dot-dashed curve in fig. 1 qualitatively appears to fit the data slightly better.

To be more quantitative, we perform a least squares fit for the compositeness scale Λ. First we multiply the differential cross section in each CDF bin by the integrated luminosity and bin-width to convert into number of events:

$$N = 4200\,\mathrm{nb}^{-1} \times \Big(\Delta E_T \int d\eta \frac{d^2\sigma}{d\eta dE_T}\Big)\mathrm{nb}. \tag{17}$$

We then examine the transformed error bars to determine the statistics obeyed by the binned events. The statistical uncertainties in each bin with transverse energy $E_T <$ 115 GeV are significantly greater than $\sqrt{N}$. We therefore exclude these non-Gaussian points from our least squares analysis. Bins in the intermediate energy range 115 GeV $<$ $E_T <$ 300 GeV contain large numbers of events and follow Gaussian statistics. At the highest transverse energies, the data bins have fewer than 25 events and are described by Poisson statistics. We assign Gaussian error bars to these last points following the discussion in ref. [19] and then treat them like Gaussian bins containing greater numbers of events.

In addition to the statistical fluctuations associated with each bin, we need to consider the systematic errors. Normalizing the theoretical number of events by the factor n removes a large, correlated systematic uncertainty. We therefore subtract an averaged percentage uncertainty from all the bins' error bars. The corrected systematic errors are then small for bins with transverse energies above 80 GeV.

We adopt the χ^2 function

$$\chi^2 = \sum_{i,j} \Delta_i (V^{-1})_{ij} \Delta_j \tag{18}$$

where $\Delta_i = N_i^{\mathrm{th}} - N_i^{\mathrm{exp}}$ represents the difference between the theoretically expected and experimentally measured number of events in the ith bin, while V denotes the covariance matrix. The statistical and systematic uncertainties for each bin are summed together in quadrature to form the diagonal entries $\sigma_{ii}^2 = \sigma_i^2(\mathrm{stat}) + \sigma_{\mathrm{i}}^2(\mathrm{sys})$ in V, while the off-diagonal elements which take into account residual bin-to-bin correlations among the systematic errors are given by $\sigma_{ij}^2 = \sigma_i(\mathrm{sys})\sigma_{\mathrm{j}}(\mathrm{sys})$ [20]. We illustrate in fig. 2 the dependence of χ^2 for 22 degrees of freedom upon Λ^{-2} over the domain -0.35 TeV^{-2} $\leq \Lambda^{-2} \leq 0.45$ TeV^{-2}

using the MT set SL structure function. [4] Various limits may readily be extracted from the clean parabola appearing in this plot. For instance, we find

$$\Lambda^{-2} = 0.113 \pm 0.080 \ \text{TeV}^{-2} \tag{19}$$

by locating the $\chi^2_{\text{min}} + 1$ points on the parabola. This translates into the asymmetrical 1σ interval

$$\Lambda = 2.98^{+2.59}_{-0.69} \ \text{TeV} \tag{20}$$

for the gluon compositeness scale. Alternatively, we may quote the more conservative lower bound

$$\Lambda > 2.03 \ \text{TeV at } 95\% \ \text{CL}. \tag{21}$$

Analogous results from the other leading and next-to-leading order distribution functions evaluated at the renormalization scales $Q^2 = E_T^2/2$ and $Q^2 = E_T^2$ are displayed in table 1. We see from the 95% lower limit entries in the last column of this table that the bound in (21) represents a conservative estimate for Λ. It compares favorably with the CDF limit for the compositeness scale associated with quark substructure.

It should soon be possible to substantially improve our limit on new gluon sector physics. The current 1992-93 Tevatron run is expected to collect a data sample five times larger than the one used in this analysis. Cross sections at higher transverse jet energies will be probed, and sensitivity to any nonrenormalizable operators in the effective Lagrangian will be enhanced. In the next few years, an integrated luminosity of 60 pb^{-1} is projected. We therefore look forward to updating our findings as new data comes forth from Batavia.

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[4] Negative values for Λ^{-2} do not imply imaginary values for Λ. Instead the phase of Λ^{-2} is simply absorbed into the dimensionless C_2 coefficient which multiplies the gluon operator O_2 in our effective Lagrangian.

Distribution function	Q^2	normalization factor	$\chi^2_{\rm min}$	Λ^{-2} / TeV^{-2}	Λ_{95} / TeV
MT	$E_T^2/2$	1.35 ± 0.01	11.76	0.113 ± 0.080	2.03
set SL	E_T^2	1.61 ± 0.01	12.17	0.081 ± 0.067	2.29
MT	$E_T^2/2$	1.17 ± 0.01	11.10	0.104 ± 0.085	2.03
set S	E_T^2	1.39 ± 0.01	11.34	0.081 ± 0.071	2.25
MT	$E_T^2/2$	1.15 ± 0.01	11.15	0.104 ± 0.085	2.03
set B1	E_T^2	1.37 ± 0.010	11.45	0.081 ± 0.072	2.24
HMRS	$E_T^2/2$	1.08 ± 0.01	11.52	0.081 ± 0.092	2.08
set B	E_T^2	1.28 ± 0.01	11.48	0.077 ± 0.077	2.22
CTEQ	$E_T^2/2$	1.25 ± 0.01	15.82	-0.131 ± 0.070	2.02
set L	E_T^2	1.48 ± 0.01	15.97	-0.117 ± 0.067	2.10
CTEQ	$E_T^2/2$	1.21 ± 0.01	12.18	0.014 ± 0.084	2.45
set MS	E_T^2	1.43 ± 0.01	12.22	0.018 ± 0.071	2.65

Table 1

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Figure Captions

Fig. 1. Inclusive jet cross section plotted against transverse jet energy E_T. The data points are the experimental measurements reported by CDF. The solid curve represents the predictions of pure QCD with no composite interactions, while the dot-dashed curve illustrates the effect of gluon operator O_2 with $\Lambda = 2$ TeV and $C_2(\Lambda) = -4\pi$. The theoretical results are based upon the leading order MT set SL distribution function evaluated at $Q^2 = E_T^2/2$.

Fig. 2. χ^2 for 22 degrees of freedom plotted as a function of Λ^{-2}.